

Sterile neutrino dark matter from right-handed neutrino oscillations

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We propose a scenario where sterile neutrino (either warm or cold) dark matter (DM) is produced through (non-resonant) oscillations among right-handed neutrinos (RHNs) and can constitute the whole DM in the Universe. We study this production mechanism in a simple setup with three RHNs, where the lightest RHN can be sterile neutrino DM whose mixing with left-handed neutrinos is sufficiently small while heavier RHNs can have non-negligible mixings with left-handed neutrinos to explain the neutrino masses by the seesaw mechanism. We also demonstrate that, in our scenario, the production of sterile RHN DM from the decay of a heavier RHN is subdominant compared with the RHN oscillation production due to the X-ray and small scale structure constraints.

I. INTRODUCTION

While it has been established that neutrinos are massive due to the discovery of neutrino oscillations [1, 2], their precise properties, such as their complete mixing parameters and their being Dirac or Majorana, however are still under active investigation. An analogous (and even more perplexing) story applies to DM whose nature remains unknown despite the ever-growing evidence for its existence from the astrophysical observables. An intriguing possibility regarding these mysteries would be to introduce RHNs which can address the origin of neutrino masses and act as DM, and their importance can well go beyond the DM and neutrino physics including their potential roles in the inflation and baryon asymmetry production [3–8].

In this letter, we seek a possibility for a sterile RHN to make up the whole DM in the Universe and, in particular, propose the new production mechanism of sterile RHN DM through the mixing among RHNs. Our production mechanism differs from the conventional active-sterile neutrino oscillation production where sterile RHN DM is produced due to its mixing with left-hand neutrinos. Those production mechanisms requiring the sterile RHN DM to couple to left-handed neutrinos are known to suffer from the severe tension between the upper DM mass bound from the X-ray data and the lower mass bound from the small-scale structure data [9–13].

These astrophysical constraints on the sterile RHN DM heavily depend on their production mechanisms and many possibilities have been explored to produce the desired sterile RHN DM abundance in addition to the conventional non-resonant/resonant active-sterile neutrino conversion mechanisms [14–17]. The examples include the RHN production from heavier particle decays, by the freeze-in and by the freeze-out accompanied by the entropy dilution, and these alternative production mechanisms typically involve the additional fields besides RHNs.

Our scenario does not introduce any additional fields besides RHNs which naturally show up in a simple extension of the SM to account for the neutrino masses. The

tight X-ray and DM lifetime bounds on the DM mass can be evaded because our scenario does not necessarily require DM mixing with SM neutrinos and the desirable sterile RHN DM abundance can be realized for both warm and cold DM mass ranges.

II. SETUP

The Lagrangian we study is the SM with three Majorana RHNs, given by

$$\begin{aligned}\mathcal{L} &= \mathcal{L}_{\text{SM}} + \mathcal{L}_N, \\ \mathcal{L}_N &= \bar{\nu}_R i \not{\partial} \nu_R - \left[\nu_R^c{}^T y_\nu L H - \frac{1}{2} \nu_R^c{}^T \mathcal{M}_N \nu_R^c + h.c. \right],\end{aligned}\tag{1}$$

where \mathcal{L}_{SM} is the SM Lagrangian, and H, L, ν_R are, respectively, Higgs doublet, lepton doublet and RHN. While we have omitted flavor indexes, the neutrino Yukawa coupling y_ν and the Majorana mass \mathcal{M}_N are understood as 3×3 matrix.

We begin with the field basis where y_ν is diagonal, denoted as y_ν^{diag} , while \mathcal{M}_N is in general a non-diagonal matrix, which we call the interaction basis in the following discussion.¹ A familiar seesaw mechanism for the mass of left-handed neutrino ν_L reads, in terms of its Dirac mass $m_D^{\text{diag}} = y_\nu^{\text{diag}} v$ (v is the Higgs VEV),

$$\mathcal{M}_\nu = m_D^{\text{diag}T} \mathcal{M}_N^{-1} m_D^{\text{diag}}\tag{3}$$

which can be diagonalized as $\mathcal{M}_\nu^{\text{diag}} = U_L^T \mathcal{M}_\nu U_L$, with U_L being the Pontecorvo-Maki-Nakagawa-Sakata matrix.² The neutrino mass eigenstates are given by

$$\begin{bmatrix} \nu_L \\ \nu_R^c \end{bmatrix} = U \begin{bmatrix} \nu \\ N^c \end{bmatrix}, \quad U = \begin{bmatrix} 1 & \theta^\dagger \\ -\theta & 1 \end{bmatrix} \begin{bmatrix} U_L \\ U_R^* \end{bmatrix},\tag{4}$$

¹ Note that the interaction basis in this letter is not the electroweak eigenstate.

² Throughout this letter, we take the charged lepton Yukawa coupling to be diagonal.

where $\theta \equiv \mathcal{M}_N^{-1} m_D^{\text{diag}}$ and U_R is a unitary matrix defined to diagonalize \mathcal{M}_N as $\mathcal{M}_N^{\text{diag}} = U_R^\dagger \mathcal{M}_N U_R^*$. By taking the rotation of Eq. (4), the Yukawa coupling y_ν is in general a non-diagonal matrix while the neutrino masses, \mathcal{M}_ν and \mathcal{M}_N , are simultaneously diagonalized. We call this field basis the mass basis. Thus, we obtain the relation

$$y_\nu^{\text{diag}} y_\nu^{\text{diag}\dagger} = v^{-2} U_R (\mathcal{M}_N^{\text{diag}})^{1/2} \mathcal{M}_\nu^{\text{diag}} (\mathcal{M}_N^{\text{diag}})^{1/2} U_R^\dagger \quad (5)$$

For other possible parametrizations, see, e.g. Refs. [18–21]. Our parametrization given by Eq. (5) is convenient for our following discussions on the estimation of the abundance and the decay rate of the RHN DM for a given RHN mass matrix. The mixing between ν_L and N is then parametrized by $\Theta = \theta^\dagger U_R^*$, and

$$\Theta^2 \equiv \Theta^\dagger \Theta = (\mathcal{M}_N^{\text{diag}})^{-1/2} \mathcal{M}_\nu^{\text{diag}} (\mathcal{M}_N^{\text{diag}})^{-1/2}. \quad (6)$$

The oscillations among RHNs can take place when their mass and interaction bases differ. We, in the following discussions, consider three RHNs with their masses $\mathcal{M}_N^{\text{diag}} = \text{diag}\{M_1, M_2, M_3\}$ and take N_1 as the lightest one so that it can play a role of DM. For the active neutrino masses, we parametrize $\mathcal{M}_\nu^{\text{diag}} = \text{diag}\{m_1, m_2, m_3\}$ for the Normal Hierarchy (NH), where $\Delta m_{21}^2 \equiv m_2^2 - m_1^2 = (7.50^{+0.19}_{-0.17}) \times 10^{-5} \text{ eV}^2$, $\Delta m_{31}^2 \equiv m_3^2 - m_1^2 = (2.457^{+0.047}_{-0.047}) \times 10^{-3} \text{ eV}^2$ [22]. For the Inverted Hierarchy (IH), we take $\mathcal{M}_\nu^{\text{diag}} = \text{diag}\{m_3, m_1, m_2\}$ and $\Delta m_{32}^2 \equiv m_3^2 - m_2^2 = (-2.449^{+0.048}_{-0.047}) \times 10^{-3} \text{ eV}^2$. The lightest neutrino mass (m_1 for the NH case, and m_3 for the IH case) is taken as a free parameter. In our discussions below, whenever it is not necessary to distinguish the mass orderings, m_1 refers to the lightest mass for brevity. We can see from Eq. (6) that Θ^2 is diagonal, so that one just needs to check Θ_{11} is sufficiently small to evade the constraints from the decay rate of N_1 such as the X-ray and lifetime bounds. A few concrete RHN mass matrices which can realize the desired DM relic abundance and neutrino masses via the seesaw mechanism are given in §IV.

III. DM PRODUCTION THROUGH RHN OSCILLATION

We now check if the enough abundance of ν_{R1} can be produced from the RHN oscillations. In our scenario, ν_{R2}, ν_{R3} explain the left-handed neutrino masses by the seesaw mechanism and they can have sizable neutrino Yukawa couplings to be in the thermal equilibrium at a sufficiently high temperature. ν_{R1} , on the other hand, has a sufficiently small coupling to the SM species, so that its production is dominated by the conversion from heavier RHNs. For the clarity of the following quantitative discussion, we focus on the ν_{R1} abundance produced only from its mixing with ν_{R2} because ν_{R3} plays the same role as ν_{R2} in producing ν_{R1} .

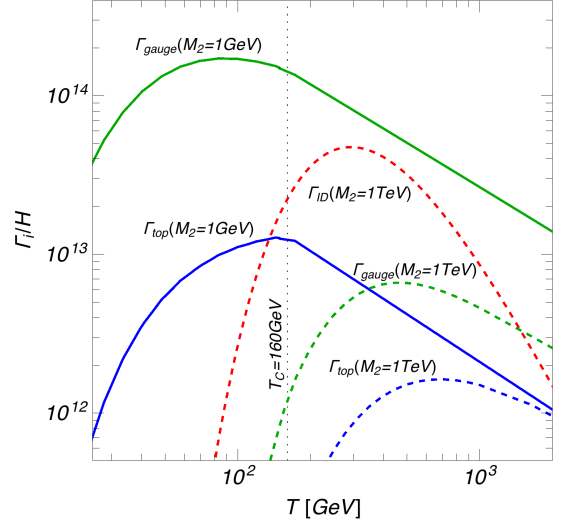


FIG. 1. The ratios between the rescaled (i.e. divided by the Yukawa couplings) reaction rates and the Hubble parameter are shown (the actual reactions rates are obtained by multiplying the Yukawa couplings). The solid curves are for $M_2 = 1 \text{ GeV}$ and the dashed curves are for $M_2 = 1 \text{ TeV}$.

The relevant reactions for the ν_{R2} thermalization are: the scatterings caused by Yukawa interaction, $\nu_{R2} L \leftrightarrow t Q_3$, $\nu_{R2} t \leftrightarrow L Q_3$, $\nu_{R2} Q_3 \leftrightarrow L t$, those involving the gauge bosons, $\nu_{R2} V \leftrightarrow H L$, $\nu_{R2} L \leftrightarrow H V$, $\nu_{R2} H \leftrightarrow L V$ and the decay and inverse decay $\nu_{R2} \leftrightarrow L H$ ($Q_3(t)$ is the left (right) handed top quark, and V represents the $SU(2)_L$ and $U(1)_Y$ gauge bosons).

The Boltzmann equation for ν_{R1} [23] reads

$$\frac{dn_{\nu_{R1}}}{dt} + 3Hn_{\nu_{R1}} = C_{\nu_{R1}} \quad (7)$$

where $C_{\nu_{R1}}$ represents the collision term integrated over the ν_{R1} momentum given by

$$C_{\nu_{R1}} \simeq \mathcal{P}(\nu_{R2} \rightarrow \nu_{R1})(\gamma_{\nu_{R2}}^{\text{col}} + \gamma_{\nu_{R2}}^{\text{ID}}), \quad (8)$$

$$\mathcal{P}(\nu_{R2} \rightarrow \nu_{R1}) = \frac{1}{2} \sin^2 2\theta_N, \quad (9)$$

$$\gamma_{\nu_{R2}}^{\text{col}} = \frac{T}{64\pi^4} \int_{s_{\min}}^{\infty} ds \hat{\sigma} \sqrt{s} K_1(\sqrt{s}/T), \quad (10)$$

$$\gamma_{\nu_{R2}}^{\text{ID}} = \frac{M_2^2 T}{\pi^2} \Gamma(\nu_{R2} \rightarrow L H) K_1(M_2/T). \quad (11)$$

Here \mathcal{P} is the oscillation probability (θ_N is the mixing angle between ν_{R1} and ν_{R2}), $\Gamma(\nu_{R2} \rightarrow L H) \simeq (y_\nu y_\nu^\dagger)_{22} M_2 / (8\pi)$ is the decay width, and $\hat{\sigma}$ is the reduced cross section for the ν_{R2} collisions with the kinematical cut s_{\min} of the Mandelstam variable s , and K_1 is the modified Bessel function of the first kind.³

³ A factor 1/2 in \mathcal{P} comes from averaging out the RHN oscillation because the oscillation time scale is much shorter

ν_{R1} is efficiently produced when the collision terms ($\propto \gamma_{\nu R2}$) are large.⁴ Fig. 1 shows Γ_i/H where Γ_i represents the rescaled reaction rates for the process i by taking the neutrino Yukawa coupling as unity (so that the curves can be easily scaled by multiplying the Yukawa coupling of interest). We define Γ_{top} and Γ_{gauge} by $\gamma_{\nu R2}^{\text{col}}/n_t^{\text{eq}}$ and $\gamma_{\nu R2}^{\text{col}}/n_V^{\text{eq}}$, respectively, where the former (latter) includes processes involving top quarks (gauge bosons), and the inverse decay rate $\Gamma_{\text{ID}} = \gamma_{\nu R2}^{\text{ID}}/n_H^{\text{eq}}$. n_i^{eq} ($i = t, V, H$) represents the equilibrium number density of species i . The figure shows the plots for $M_2 = 1$ GeV (solid) and for $M_2 = 1$ TeV (dashed), and we note that the inverse decay takes place only for the latter because of the kinematics. The actual reaction rates can be obtained by multiplying these rescaled reaction rates by $(y_\nu y_\nu^\dagger)_{22}$. We can see, from Fig. 1, that N_2 can reach the thermal equilibrium ($\Gamma_i/H \gtrsim 1$) when $(y_\nu y_\nu^\dagger)_{22}$ is larger than $10^{-14} - 10^{-13}$ for $M_2 = 1 - 10^3$ GeV, which is also in the desired numerical range to explain the neutrino masses by the seesaw mechanism.

The produced ν_{R1} (interaction state) constitutes the DM N_1 (mass eigenstate),⁵ and the current N_1 relic number density can be estimated, in terms of the yield parameter $Y_{N_1} \equiv n_{N_1}/s$ (s is the entropy density), by integrating the Boltzmann equation from T_{RH} , the reheating temperature, to the current temperature $T = T_0$

$$Y_{N_1}^0 \equiv Y_{N_1}(T=0) = \int_0^\infty dT \mathcal{P}(\nu_{R2} \rightarrow \nu_{R1}) \frac{\gamma_{\nu R2}}{sHT} \quad (12)$$

where we have taken the limits $T_{\text{RH}} \rightarrow \infty, T_0 \rightarrow 0$, and $\gamma_{\nu R2} \equiv \gamma_{\nu R2}^{\text{col}} + \gamma_{\nu R2}^{\text{ID}}$. The corresponding DM density can

than the collision time scale involving ν_{R2} . More quantitatively, this averaging is justified for $T \lesssim 10^6$ GeV and/or $\Delta M^2 \equiv M_2^2 - M_1^2 \gtrsim 1$ GeV² because $t_{\text{osc}}/t_{\text{col}} \sim (y_\nu^2/10^{-14})(g^2/10^{-2})(\text{GeV}^2/\Delta M^2)(T/10^6 \text{ GeV})^2$ where g represents a gauge coupling for a relevant gauge interaction. As we will discuss later, y_ν^2 of order 10^{-14} is required for GeV-scale RHN to reach the thermal equilibrium and it is automatically realized by enforcing the seesaw mechanism. The finite temperature effects on the RHN mixing angle θ_N are suppressed by the neutrino Yukawa couplings in our scenario and we simply consider a constant θ_N in our estimation. The cases when these approximations are not applicable are left for the future work.

⁴ Some of collision terms, such as $\nu_R H \rightarrow LV$, possess the infrared divergences, which are regulated by the thermal mass of the propagator in our analysis for $T > T_C$ (T_C is the critical temperature of the electroweak phase transition and we take $T_C = 160$ GeV) [24–26].

⁵ The produced ν_{R1} is composed of N_1 and N_2 which propagate with different velocities. As the ν_{R1} energy gets redshifted, these two mass states are eventually well separated and thus ν_{R1} is expected to mostly develop the N_1 component as long as $M_1 \ll M_2$, although the oscillation property may call for a careful study [27].

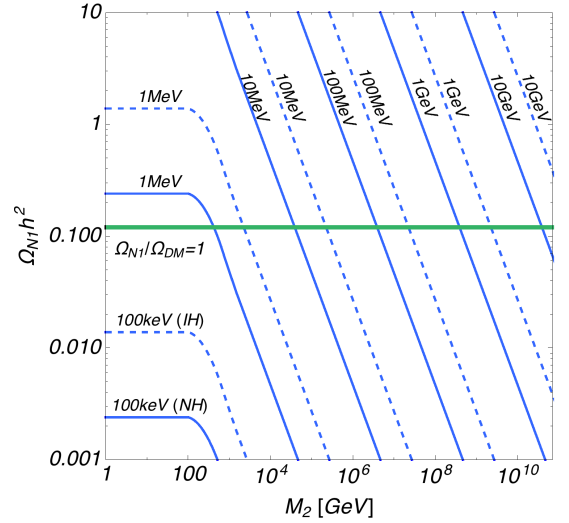


FIG. 2. The N_1 relic abundance is shown as a function of M_2 by varying M_1 from 100 keV to 10 GeV. The solid and dashed curves show the NH and IH cases, respectively.

then be estimated in terms of the yield parameter

$$\Omega_{N_1} h^2 \simeq 0.12 \left[\frac{\sin^2 2\theta_N}{8.8 \times 10^{-3}} \right] \left[\frac{(y_\nu y_\nu^\dagger)_{22}}{10^{-14}} \right] \left[\frac{M_1}{\text{keV}} \right] \left[\frac{\tilde{Y}_{N_1}^0}{10^{13}} \right], \quad (13)$$

where $\tilde{Y}_{N_1}^0$ is the rescaled yield parameter, defined by factoring out the oscillation probability and the Yukawa coupling, $\tilde{Y}_{N_1}^0 \equiv Y_{N_1}^0/(\mathcal{P}(\nu_{R2} \rightarrow \nu_{R1})(y_\nu y_\nu^\dagger)_{22})$. We found the following simple fitting formula to grasp the characteristic features of the DM abundance in our scenario

$$\begin{aligned} \log_{10} \tilde{Y}_{N_1}^0 &\simeq 13 \quad (M_2 \lesssim 100 \text{ GeV}) \\ &\simeq 13 - \log_{10}(M_2/100 \text{ GeV}) \quad (M_2 \gtrsim 100 \text{ GeV}). \end{aligned} \quad (14)$$

This behavior matches our expectation because, as emphasized in referring to Fig. 1, the most efficient production occurs when the production rate reaches maximal with respect to the Hubble expansion rate. $\tilde{Y}_{N_1}^0$ is hence little dependent on M_2 when M_2 is smaller than $T \sim T_c$, because the SM particles obtain the masses from the Higgs VEV and the most efficient production occurs around the critical temperature regardless of M_2 due to the Boltzmann suppressed production rate after the electroweak phase transition. For $M_2 \gtrsim 100$ GeV, on the other hand, the SM particles possess the thermal mass and the production rate becomes maximal around $T \sim M_2$, which leads to some power dependence of the yield parameter on M_2 . This is illustrated through a concrete example in the next section.

IV. BENCHMARK MODELS

We here provide two simple RHN mass matrices to exemplify the previous discussion for the DM abundance estimation, starting with a simple matrix

$$\mathcal{M}_N = \begin{bmatrix} m & & \\ m & M_2 & \\ & & M_3 \end{bmatrix}, \quad (15)$$

where $m \ll M_2, M_3$. \mathcal{M}_N is then diagonalized as $\mathcal{M}_N^{\text{diag}} = \text{diag}\{-m^2/M_2, M_2, M_3\}$ by using U_R which reads

$$U_R^* \simeq \begin{bmatrix} 1 & \theta_N & \\ -\theta_N & 1 & \\ & & 1 \end{bmatrix} \quad (16)$$

with $\theta_N = m/M_2$. Since N_1 mass is $M_1 = m^2/M_2$, $\theta_N = (M_1/M_2)^{1/2}$.

For the NH case, the diagonalized Yukawa coupling is given by $(y_\nu^{\text{diag}} y_\nu^{\text{diag}\dagger})_{22} = m_2 M_2/v^2$, and we thus obtain

$$\Omega_{N_1} h^2 \simeq 0.12 \left[\frac{m_2}{0.01 \text{ eV}} \right] \left[\frac{M_1}{0.52 \text{ MeV}} \right]^2 \left[\frac{\tilde{Y}_{N_1}^0}{10^{13}} \right]. \quad (17)$$

We note that in this case the lightest neutrino mass should be vanishingly small, $\Theta_{11}^2 = m_1/M_1 \approx 0$ (which can even vanish without affecting our scenario), so that the X-ray/lifetime constraints can be avoided, and $m_2 \simeq \sqrt{\Delta m_{21}^2}$. In the IH case, m_2 in Eq. (17) should be replaced by $m_1 \simeq \sqrt{|\Delta m_{32}^2|}$ by demanding $\Theta_{11}^2 = m_3/M_1 \approx 0$. In this example, Fig. 2 shows $\Omega_{N_1} h^2$ as a function of M_2 for various M_1 taken from 100 keV to 10 GeV in both the NH and IH cases which are depicted by solid and dashed curves, respectively.⁶ The green band in the figure indicates the observed value of the DM abundance given by $\Omega_{\text{DM}} h^2 = 0.1197 \pm 0.0022$ [28].

Next, let us consider another example where \mathcal{M}_N is given by

$$\mathcal{M}_N = \begin{bmatrix} M_0 & m & \\ m & M_2 & \\ & & M_3 \end{bmatrix}, \quad (18)$$

and m and M_0 are taken to be $m, M_0 \ll M_2, M_3$. By using the same U_R as given in Eq. (16), we obtain $\mathcal{M}_N^{\text{diag}} = \text{diag}\{M_0 - m^2/M_2, M_2, M_3\}$. With $\theta_N = m/M_2$, the resultant N_1 abundance in the NH case is given by

$$\Omega_{N_1} h^2 \simeq 0.12 \left[\frac{m_2}{0.01 \text{ eV}} \right] \left[\frac{M_1}{\text{keV}} \right] \left[\frac{(m/5 \text{ GeV})^2}{M_2/100 \text{ GeV}} \right] \left[\frac{\tilde{Y}_{N_1}^0}{10^{13}} \right], \quad (19)$$

while, in the IH case, m_2 should be replaced by m_1 . The DM relic abundance is independent of the mixing angle $\Theta_{11}^2 (= m_1/M_1 (m_3/M_1)$ for the NH(IH)), and the X-ray and lifetime constraints can be evaded as in the previous example.

⁶ It should be noted that, in Fig. 2, $t_{\text{osc}}/t_{\text{col}} \ll 1$ is achieved for $T \lesssim 10^6 \times M_2$ even in the large M_2 region, so that a factor $1/2$ in \mathcal{P} by averaging out the RHN oscillation is justified.

V. DISCUSSION/CONCLUSION

Before concluding our discussions, let us briefly point out another potentially interesting production mechanism: the production of N_1 from a heavier RHN decay. We can consider the decay of N_2 (and/or N_3) which is thermally decoupled while it is relativistic (otherwise N_2 number density would be too small due to the Boltzmann suppression). N_1 abundance then can be estimated as

$$\Omega_{N_1} h^2 \simeq 10^{-10} \left[\frac{\Theta_{11}^2}{10^{-12}} \right] \left[\frac{M_1}{10 \text{ keV}} \right] \left[\frac{g_*(T_0)}{g_*(T_{\text{FO}})} \right] \quad (20)$$

where we used the branching fraction of N_2 decay for the process $N_2 \rightarrow N_1 + (\text{mesons, leptons})$, $\text{Br}(N_2 \rightarrow N_1) \simeq \Gamma(N_2 \rightarrow N_1)/\Gamma(N_2 \rightarrow SM) \simeq M_2 \Theta_{11}^2 \Theta_{22}^2 / M_2 \Theta_{22}^2 \simeq \Theta_{11}^2$, and the ratio of g_* accounts for the change in the effective degrees of freedom from the N_2 freeze-out epoch to the present time. This production contribution is hence subdominant compared with RHN oscillation production in the parameter region of our interest.

Let us next mention the small scale structure constraints applicable to our scenario. We here discuss the Lyman- α forest constraints which can give the lower limit on the DM mass from the DM free streaming scale $\lambda_{FS} \sim 1 \text{ Mpc}(\text{keV}/M_1)(\langle p/T \rangle/3.15)$ [29]. Too large a free streaming scale can be excluded due to the suppression of small-scale structure formation. The average momentum of N_1 produced by the non-resonant oscillation of thermalized N_2 can be estimated as $\langle p_1 \rangle \sim 2.8T$, analogous to the conventional (non-resonant) active-sterile oscillation scenario. Taking account of momentum redshifting by a factor $(g(T_{N_2 \rightarrow N_1})/g(T \ll \text{MeV}))^{-1/3}$ due to the change in the effective degrees of freedom, Lyman- α data leads to the RHN DM mass bound $M_1 \gtrsim 10 \text{ keV}$ for our scenario [12] (when $N_2 \rightarrow N_1$ occurs most efficiently before the QCD phase transition which is the case for the parameter range discussed so far). Such a DM mass range can be realized in our scenario as explicitly demonstrated through the concrete examples in the last section while being compatible with both the right relic abundance and seesaw mechanism.

We have proposed the sterile neutrino DM production mechanism through the mixing among RHNs. Our scenario does not suffer from the X-ray constraints and the RHN DM mass M_1 can be above keV and even above MeV scale, so that a wide range of DM properties are possible covering both warm and cold DM features. Because of the simplicity of our proposed model in that no other fields besides RHNs are introduced, our scenario would potentially open up a new avenue to further explore the interplay between the neutrino physics and DM phenomenology.

Among the possible extensions of our DM scenarios, we plan to study the leptogenesis as well as the neutrino observables such as the neutrinoless double beta decay in our future work. For instance, even though we have focused on the DM production in this letter, the neu-

trino Yukawa couplings in our model can be further constrained by seeking the production of desirable baryon asymmetry in the Universe. The realization of leptogenesis when N_2 and N_3 are heavy enough and/or are degenerate in their masses with sufficient CP violations [5, 6, 24] will be explored in our forthcoming paper. The CP phases in the neutrino Yukawa couplings are of great importance not only for the leptogenesis but also for the DM production in our scenario, and the presented production mechanism for the RHN DM would uncover a

new connection between DM and leptogenesis to bring considerable opportunities for subsequent studies.

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